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The Contribution of Different Supernova Populations to the Galactic Gamma-Ray Background

E.G. Berezhko

Institute of Cosmophysical Research and Aeronomy, Lenin Ave. 31, 677891 Yakutsk, Russia

berezhko@ikfia.ysn.ru

and

H. J. Völk

Max–Planck–Institut für Kernphysik, P.O. Box 103980, D-69029 Heidelberg, Germany

Heinrich.Voelk@mpi-hd.mpg.de

ABSTRACT

The contribution of Source Cosmic Rays (SCRs), accelerated and still confined in Supernova Remnants (SNRs), to the diffuse high energy γ-ray emission above 1 GeV from the Galactic disk is studied. The γ-rays produced by the SCRs have a much harder spectrum than those generated by the Galactic Cosmic Rays (GCRs). Extending a previous paper, a simple SNR population synthesis is considered and the Inverse Compton emission from the SCR electrons is evaluated in greater detail. Then the combined spectrum of γ-ray emission from the Galactic Supernova Remnant population is calculated and this emission at low Galactic latitudes is compared with the diffuse γ-ray emission observed by the EGRET and ground based instruments. The average contribution of SCRs is comparable to the GCR contribution already at GeV energies, due to Supernovae of types II and Ib exploding into the wind bubbles of quite massive progenitor stars, and becomes dominant at γ-ray energies above 100 GeV. At TeV energies the dominant contribution is from SCRs in SNRs that expand into a uniform interstellar medium. In fact, the sum of hadronic and Inverse Compton γ-rays would exceed the limits given by the existing experimental data, unless the confinement time T_{SN}, i.e. the time until which SNRs confine the main fraction of accelerated SCRs, is as small as $T_{SN} \sim 10^4$ yr and the typical magnetic field strength in SNRs as large as 30 μG. Both situations are however possible as a result of field amplification through CR backreaction in the acceleration process. It is pointed out that accurate measurements of the low-latitude diffuse Galactic γ-ray spectrum at TeV-energies can serve as a unique consistency test for CR origin from the Supernova Remnant population as a whole.

Subject headings: gamma-rays – background radiation – cosmic rays – supernova remnants

1. Introduction

The diffuse Galactic γ-ray emission was measured with the Energetic Gamma Ray Experiment Telescope (EGRET) on the Compton Gamma Ray Observatory up to about 20 GeV. To first approximation it can be described by a suitable model for the diffuse interstellar gas, the Galactic Cosmic Ray (GCR) distribution, and the diffuse photon fields (Hunter et al. 1997a,b). The resulting GCR distribution in the Milky Way essentially corresponds to that measured in situ in the local neighborhood of the Solar System. This is also our starting point.

However above about 1 GeV the observed average diffuse γ-ray intensity in the inner Galaxy, $300^\circ < l < 60^\circ$, $|b| \leq 10^\circ$, exceeds the model prediction significantly, and this excess appears to increase monotonically with γ-ray energy. Spatially, the diffuse γ-ray intensity is strongly increasing towards the Galactic midplane.

In a previous paper Berezhko & Völk (2000a, hereafter referred to as Paper I), we have calculated the contribution of an unresolved distribution of CR sources to the diffuse γ-ray flux in the disk, assuming the GCR distribution within its large confinement volume to be the same as that measured in situ. The CR sources contribute through the γ-ray emission of the accelerated particles that are confined in their interior – the Source Cosmic Rays (SCRs). It was assumed that Supernova Remnants (SNRs) are the dominant sources of the GCRs. On this premise it was found that the SCRs give an important contribution to the diffuse high-energy γ-ray emission at low Galactic latitudes. Since the CR energy spectrum inside SNRs is much harder than on average in the Galaxy — the average spectrum being softened by rigidity-dependent escape from the Galaxy in the diffusion/convection region above the disk — the relative SNR contribution increases with energy and was found to become in fact dominant at γ-ray energies $\epsilon_\gamma \gtrsim 100$ GeV. This led to a substantial increase of the "diffuse" TeV γ-ray emission from the Galactic disk so as to constitute a significant observational background. This hard-spectrum background must also be taken into account in the search for spatially extended Galactic CR sources above the GeV region.

Extending Paper I we take into account here that a significant fraction of the SNR population, the type II and Ib Supernovae (SNe) with progenitor masses in excess of $15 M_\odot$, expand into a nonuniform circumstellar medium strongly modified by the wind from the progenitor star. In such SNRs the majority of CRs should be produced in a thin dense shell

of swept-up ISM. As a consequence of the high density of this shell the expected π^0-decay γ-ray emission with energies $\epsilon_\gamma \lesssim 100$ GeV is shown to be an order of magnitude larger than for an identical explosion into a uniform ISM. This gives a natural explanation for the observed diffuse γ-ray excess for energies $\gtrsim 1$ GeV. We also consider in detail the contribution of the high-energy SCR electrons to the high energy γ-ray background due to inverse Compton (IC) scattering on the background photon field. Since the energy spectrum of the SCR electrons is strongly influenced by synchrotron losses, the IC γ-ray emission depends on the effective magnetic field strength in the SNR interior. This field may in turn be amplified relative to the external field by the acceleration process itself (Bell & Lucek 2001). We take such a field amplification into account as an average property of SNRs relative to the contribution from the GCRs. As demonstrated below, this results in an order of magnitude increase of the Galactic γ-ray emission at TeV-energies. In a preliminary form these results have been presented in Berezhko & Völk (2003).

We therefore conclude that a detailed measurement of the low-latitude diffuse Galactic γ-ray spectrum within the energy interval from 1 to 10^4 GeV will constitute an important indirect test for the GCR origin from the *population* of Galactic SNRs as a whole.

2. Gamma-ray luminosity of SNRs expanding into a uniform circumstellar medium

The majority of the GCRs with charge number Z, at least up to kinetic energies $\epsilon \sim 10^{15} Z$ eV, is presumably accelerated in SNRs. Individual examples, where this is the case, are the objects SN 1006 (Berezhko et al. 2002, 2003a) and Cas A (Berezhko et al. 2003b; Berezhko & Völk 2004). According to theory a significant part of the hydrodynamic SN explosion energy $E_{SN} \approx 10^{51}$ erg is converted into CRs already in the early Sedov phase of the evolution, $t \sim 10^3$ yr, as a result of diffusive shock acceleration (Berezhko et al. 1996; Berezhko & Völk 1997, 2000b). Later on the total CR energy content E_c and therefore the γ-ray production slowly varies with time. We neglect this variation in the energetics in our considerations below, assuming a time-independent SCR energy content $E_c = 0.1 E_{SN}$, consistent with the GCR energy budget requirement.

The total number of active SNRs, $N_{SN} = \nu_{SN} T_{SN} \approx 300$ to 3000, is an increasing function of their assumed life time $T_{SN} = 10^4$ to 10^5 yr, i.e. the time until which they can confine the main fraction of the accelerated particles. Here $\nu_{SN} \approx 1/30$ yr is the Galactic SN rate. Consequently we conclude that the population of the oldest SNRs should dominate the total γ-ray luminosity of the ensemble of Galactic SNRs. Approximately we shall therefore consider only SNRs in the Sedov phase of their evolution in our estimate of the total γ-ray

luminosity.

Gamma-ray producing CRs in the Galaxy are then the sum of two basically different populations. The first one consists of the ordinary GCRs and presumably occupies a large Galactic residence volume quasi-uniformly. This residence volume exceeds by far the volume of the gas disk which harbors the CR sources (Ptuskin et al. 1997; Berezinsky et al. 1990). The second CR population are the SCRs in the localized SNRs.

During the initial, active period of SNR evolution at times $t \lesssim T_{SN}$ when the SN shock is relatively strong, the volume occupied by the accelerated CRs practically coincides with the shock volume. During later evolutionary stages the shock becomes weak and CRs begin to leave the SNR acceleration region. After some period of time the escaping SCRs become very well mixed with the ambient GCRs. Thus the controling factor is the shock strength.

The γ-ray production by the GCRs is quite well studied (Hunter et al. 1997b; Mori 1997). Therefore we shall concentrate on the relative contribution of the SCR population.

It was shown in Paper I that the total γ-ray spectrum measured from an arbitrary Galactic disk volume is expected to be

$$\frac{dF_\gamma}{d\epsilon_\gamma} = \frac{dF_\gamma^{GCR}}{d\epsilon_\gamma}[1.4 + R(\epsilon_\gamma)], \tag{1}$$

where $dF_\gamma^{GCR}/d\epsilon_\gamma$ is the π^0-decay γ-ray spectrum due to GCRs and the additional factor 0.4 is introduced to approximately take into account the contribution of GCR electron component to the diffuse γ-ray emission at GeV energies (Hunter et al. 1997b).

For the ratio $R(\epsilon_\gamma) = Q_\gamma^{SCR}/Q_\gamma^{GCR}$ of the γ-ray production rates due to SCRs and GCRs we have

$$R(\epsilon_\gamma) = 0.07\zeta \left(\frac{T_p}{10^5 \text{ yr}}\right)\left(\frac{\epsilon_\gamma}{1 \text{ GeV}}\right)^{0.6}(1 + R_{ep}), \tag{2}$$

with $\zeta = N_g^{SCR}/N_g^{GCR}$, where $N_g^{GCR} = 1cm^{-3}$ and N_g^{SCR} are the gas number density in the Galactic disk and inside SNRs, respectively, $R_{ep} = Q_\gamma^{IC}/Q_\gamma^{pp}$ is the ratio of the IC to π^0-decay γ-ray production rates due to the SCRs, T_p is the proton SCR confinement time, and the difference between the spectral indices of the GCRs and the SCRs is assumed to be 0.6.

In the case of SNRs expanding into a uniform ISM T_p, which is equal to the confinement time T_e of the SCR electron component if synchrotron losses are not more restrictive for electrons (see below), is given by the expression

$$T_p = \min\{T_{SN}, 10^3(\epsilon/\epsilon_{max})^{-5} \text{ yr}\}, \tag{3}$$

where $\epsilon_{max} = 10^5 Z$ GeV is the maximal energy of CRs that can accelerated in SNRs, and $\epsilon_\gamma = 0.1\epsilon$ which is roughly valid for the hadronic γ-ray production process. The decrease of

the proton confinement time T_p with energy ϵ is due to the diminishing ability of the SNR shock to produce high-energy CRs during the Sedov phase. In fact, the highest CR energy ϵ and therefore the highest-energy γ-ray production slowly decrease with time as $\epsilon \propto t^{-1/5}$ which implies the escape of the highest energy CRs from the SNR (Berezhko 1996; Berezhko et al. 1996; Berezhko & Völk 2000a).

Since it was shown in Paper I that electron SCRs give a γ-ray contribution comparable with that of the nuclear SCRs, we consider the electron emission here in greater detail. In the relevant energy range the γ-ray luminosity due to the electron IC scattering on the background photons can be described in the Thompson limit (Longair 1981; Berezinsky et al. 1990). Since the electron spectrum in a SNR depends significantly on the SNR age t due to synchrotron losses, we first determine the γ-ray production rate $dQ_\gamma^{IC}(\epsilon_\gamma)/dN_{SN}^e$ due to one SNR of age t:

$$\frac{dQ_\gamma^{IC}(\epsilon_\gamma)}{dN_{SN}^e} = \sigma_T c N_{ph} \frac{d\epsilon_e}{d\epsilon_\gamma} \frac{dn_{SCR}^e(\epsilon_e)}{dN_{SN}^e}, \tag{4}$$

where

$$\epsilon_e = m_e c^2 \sqrt{3\epsilon_\gamma/(4\epsilon_{ph})} \tag{5}$$

is the energy of electrons which produce an IC photon with mean energy ϵ_γ, $\sigma_T = 6.65 \times 10^{-25}$ cm^2 is the Thomson cross section, and ϵ_{ph} and N_{ph} are the mean energy and number density of the background photons, respectively. Finally

$$dn_{SCR}^e/dN_{SN}^e = N_{SCR}^e/V_g \tag{6}$$

denotes the spatial electron SCR number density averaged over the Galactic disk volume V_g, where $dN_{SN}^e = \nu_{SN} dt$ is the number of SNRs of age between t and $t + dt$ which contribute to the IC emission at energy ϵ_γ from CR electron sources.

The source electron spectrum at energies $\epsilon < \epsilon_l$, where synchrotron losses are not important, can be represented in the form

$$N_{SCR}^e(t, \epsilon) = K_{ep} N_{SCR}(\epsilon), \tag{7}$$

where $K_{ep} \approx 10^{-2}$ is the electron to proton ratio and $N_{SCR}(\epsilon)d\epsilon$ is the total number of SCR protons per SNR in the energy interval dϵ. Finally

$$\epsilon_l = 1.25 \left(\frac{10^5 \text{ yr}}{t}\right) \left(\frac{10\ \mu\text{G}}{B}\right)^2 \text{ TeV} \tag{8}$$

is the low energy limit of that part of the electron spectrum which is modified by the synchrotron losses, and B is the magnetic field strength inside the SNRs. For $\epsilon > \epsilon_l$ the electron spectrum is steeper due to synchrotron losses:

$$N_{SCR}^e(t, \epsilon) = K_{ep} N_{SCR}(\epsilon)(\epsilon_l/\epsilon). \tag{9}$$

We also neglect the time dependence of the proton spectrum in SNRs, since in the later Sedov phase which is most important here, it is not significant, especially at high energies (Berezhko et al. 1996). The proton spectrum is taken in a power lower form $N_{SCR} \propto \epsilon^{-\gamma}$ with $\gamma = 2.15$, consistent with the requirements for the sources of the Galactic CRs.

The maximum energy of SCR electrons is also restricted by their synchrotron losses during their acceleration at the SNR shock:

$$\epsilon^e_{max} = 24 \left(\frac{V_s}{10^3 \text{ km/s}}\right) \left(\frac{B}{10 \ \mu\text{G}}\right)^{-1/2} \text{ TeV}, \tag{10}$$

if $\epsilon^e_{max} < \epsilon_{max}$; here V_s is the shock speed. At the beginning of the Sedov phase ($t \sim 10^3$ yr) the shock speed is about $V_s \approx 4 \times 10^3$ km/s and subsequently decreases as $V_s \propto t^{-3/5}$. Therefore the electron confinement time can be written in the form

$$T_e = \min\{T_p, 10^3(\epsilon_e/\epsilon^e_{max})^{-5/3} \text{ yr}\}. \tag{11}$$

The total IC γ-ray production rate due to all existing SNRs can be determined by the expression

$$Q^{IC}_\gamma = \int_0^{T_e} \frac{dQ^{IC}_\gamma}{dN^e_{SN}} \nu_{SN} dt. \tag{12}$$

First we normalize Q^{IC}_γ to the production rate of π^0-decay γ-rays from inelastic CR - gas collisions, primarily p-p collisions, which may be written in the form (Drury et al. 1994)

$$Q^{pp}_\gamma(\epsilon) = Z_\gamma \sigma_{pp} c N_g n(\epsilon), \tag{13}$$

where N_g is the local gas number density, σ_{pp} is the inelastic p-p cross-section, Z_γ is the so-called spectrum-weighted moment of the inelastic cross-section, $n(\epsilon) = N_{SCR}(\epsilon)\nu_{SN} T_p/V_g$ is the spatial number density of SCRs averaged over the Galactic volume V_g, and c is the speed of light. Then, performing the integration over all possible SNR ages, we find the ratio of luminosities:

$$R_{ep} = 73.6 K_{ep} \left(\frac{N_{ph}}{N^{SCR}_g}\right) \left(\frac{4\epsilon_\gamma \epsilon_{ph}}{3 m_e^2 c^4}\right)^{(\gamma-1)/2} \tag{14}$$

for γ-ray energies $\epsilon_\gamma < \epsilon^*_\gamma$, and

$$R_{ep} = 73.6 K_{ep} \left(\frac{N_{ph} T_e}{N^{SCR}_g T_p}\right) \left(\frac{4\epsilon_\gamma \epsilon_{ph}}{3 m_e^2 c^4}\right)^{(\gamma-1)/2} \left(\frac{\epsilon_\gamma}{\epsilon^*_\gamma}\right)^{-1/2} \left(1 + \ln\sqrt{\frac{\epsilon_\gamma}{\epsilon^*_\gamma}}\right) \tag{15}$$

for $\epsilon_\gamma > \epsilon^*_\gamma$, where

$$\epsilon^*_\gamma \approx 8 \left(\frac{\epsilon_{ph}}{1 \text{ eV}}\right) \left(\frac{10^5 \text{ yr}}{T_e}\right)^2 \left(\frac{B}{10 \ \mu\text{G}}\right)^{-4} \text{ TeV} \tag{16}$$

is the energy of γ-rays which are emitted by electrons with synchrotron loss time equal to T_e. The value $Z_\gamma = 0.113$ was used, which corresponds to $\gamma = 2.15$ (Drury et al. 1994).

Due to the rather hard SCR electron spectrum the dominant contribution to the IC radiation comes from collisions with cosmic microwave background (CMB) photons which are characterized by $N_{ph} = 400$ cm^{-3} and $\epsilon_{ph} = 6.7 \times 10^{-4}$ eV. The far infrared radiation (FIR) field with $N_{ph} \approx 20$ cm^{-3} and $\epsilon_{ph} \approx 0.01$ eV contributes about 20% for $\epsilon_\gamma \gtrsim 100$ GeV. In the case of the FIR background the Klein-Nishina cross-section must be used instead of the Thomson limit at TeV γ-ray energies. We take account of the Klein-Nishina decrease of the cross-section at high energies through multiplying σ_T by the correction factor

$$f_{KN} = \exp(-3.5\sqrt{\epsilon_\gamma \epsilon_{ph}/m_e^2 c^4}), \tag{17}$$

which approximately describes the reduction effect (Blumenthal & Gould 1970). For $\epsilon_\gamma =$ 10 TeV, $f_{KN} \approx 0.1$ for the FIR contribution.

Bremsstrahlung γ-rays play no role for the average γ-ray background above GeV energies, if $K_{ep} \ll 0.1$ (Paper I).

In Fig.1 we present a calculated γ-ray spectrum based on the above expressions (1)–(17) with $T_{SN} = 10^5$ yr, $B = 10$ μG, $\epsilon_{max} = 10^5$ GeV, $K_{ep} = 10^{-2}$ and $N_g^{SCR} = N_g^{GCR}$. A SCR power law index $\gamma = 2.15$ and a 10 percent efficiency of SCR production in SNR, required to account for the GCRs, were used.

One can see that the SCR contribution becomes dominant for $\epsilon_\gamma \gtrsim 100$ GeV where the expected γ-ray spectrum becomes extremely hard. At TeV energies the predicted flux exceeds the lowest HEGRA upper limit (Aharonian et al. 2001) almost by a factor of two. The SCR electron contribution increases with energy ϵ_γ and exceeds the proton SCR contribution for energies $\epsilon_\gamma \gtrsim 300$ GeV. At $\epsilon_\gamma \sim 1$ TeV electron SCRs contribute about 60% of the total γ-ray flux. Therefore at these high energies the expected γ-ray emissivity of SNRs is only weakly dependent on the most relevant parameter, the maximum confinement time T_{SN}, since according to expressions (2) and (15) we have for $\epsilon_\gamma \sim 1$ TeV

$$dF_\gamma^{IC}/d\epsilon_\gamma \propto 1 + 0.3\ln(T_{SN}/10^5 \text{ yr}), \tag{18}$$

taking into account that at this energy $T_p = T_e = T_{SN}$. Even if T_{SN} is as short as 3×10^4 yr, the expected SCR contribution at TeV energies is therefore only 30% lower and still exceeds the GCR contribution by more than an order of magnitude (see Fig.1). We emphasize that in the latter case the SCR contribution is entirely due to the electrons: they radiate about 85% of the TeV γ-rays : This extreme case has to be considered as the lowest value for the SCR contribution to the Galactic γ-ray background radiation, since there are no other ways

to decrease it. This holds in particular also for the average gas number density N_g^{SCR} inside SNRs because the IC γ-ray production rate does not depend on it. Such a short confinement time $T_{SN} \sim 10^4$ yr (Ptuskin & Zirakashvili 2003) could be realized in a scenario where a strong amplification of the magnetic field B in SNRs is balanced by nonlinear wave-wave interactions such that the wave magnetic power $P(l)$ in the spatial scale l is dissipated in a (minimal) eddy turnover time $\sim l/v(l)$ (Verma et al. 1996). Here $v(l) = v_A[P(l)/B^2]^{1/2}$ is the rms-value of the plasma mass velocity at scale l, and v_A denotes the Alfvén velocity.

We also point out that the restriction of the electron confinement time T_e due to their synchrotron losses becomes significant for γ-ray energies $\epsilon_\gamma > 1$ TeV and leads to the steepening of the γ-ray spectrum (see Fig.1).

According to Fig.1, for energies $\epsilon_\gamma = 0.4$ to 4 TeV the expected differential γ-ray energy spectrum from the central part of the Galaxy has the form

$$dF_\gamma/d\epsilon_\gamma = A(\epsilon_\gamma/1 \text{ TeV})^{-\alpha} \qquad (19)$$

with power law index $\alpha \approx 2.25$ and amplitude $A = (6 \ to \ 8) \times 10^{-9}$ (cm^2 s sr TeV)$^{-1}$, depending on the confinement time T_{SN}.

As one can see from Fig.1, the γ-ray flux expected for $T_{SN} = 10^5$ yr at $\epsilon_\gamma = 1$ TeV exceeds the HEGRA upper limit (Aharonian et al. 2001) and the preliminary flux measured by the Milagro detector (Fleysher 2003). We emphasize here that we have multiplied the Milagro flux by a factor of 3 since we present the expected γ-ray flux from the region $|b| \leq 2°$ whereas the Milagro data correspond to $|b| \leq 5°$. Even for $T_{SN} = 3 \times 10^4$ yr the expected flux exceeds the Milagro data.

The only physical parameter which strongly influences the IC γ-ray production rate is the SNR magnetic field B: according to Eqs. (15)–(16) $F_\gamma^{IC} \propto B^{-2}$. A value of B that is significantly higher than $B = 10\ \mu$G can be attributed to field amplification at the shock front due to the strong wave production by the acceleration of CRs far into the nonlinear regime (Bell & Lucek 2001).

In Fig.2 we present the same calculations as in Fig.1 but with a mean SNR magnetic field value $B = 30\ \mu$G. Compared with the previous case the IC γ-ray emission is decreased by an order of magnitude. Therefore the contribution of the hadronic SCRs to the γ-ray flux becomes much more significant: at TeV-energies and for $T_{SN} = 10^5$ yr the π^0-decay γ-rays exceed the IC γ-rays by a factor of 2.5 , whereas for $T_{SN} = 3 \times 10^4$ yr the IC γ-rays still exceed the π^0-decay γ-rays by a factor of 1.3. One can see that in this case the expected γ-ray flux is below the HEGRA and Tibet upper limits already for a SCR confinement time $T_{SN} = 10^5$ yr. The expected γ-ray spectrum at energies $\epsilon_\gamma = 0.1 \ to \ 10$ TeV is characterized by $\alpha = 2.4$ and $A = (2.4 \ to \ 4) \times 10^{-9}$ (cm^2 s sr TeV)$^{-1}$, which also considerably exceeds

the GCR contribution. Note that since the maximum energy of SCRs obeys $\epsilon_{max} \propto B$, it is larger by a factor of three compared with the previous case. Taking into account the uncertainty of the Milagro flux one can conclude that this measurement does not contradict our prediction for the case of the amplified magnetic field. The lower value of the confinement time $T_{SN} \sim 10^4$ yr might be preferable in this sense.

3. Wind SNRs of type II and Ib

The above consideration corresponds to the assumption that all existing Galactic SNRs expand into a uniform circumstellar medium with a density of $N_g^{GCR} = 1$ cm^{-3}. However, for stellar masses in excess of about 15 $M_\odot$ the progenitors of type II and Ib SNe can strongly modify their environment through their ionizing radiation and very energetic winds (Weaver et al. 1977; Chevalier & Liang 1989). Disregarding large-scale turbulent mixing processes in the sequel, the stellar wind from such a star creates a low-density bubble beyond the termination shock out to a radius R_{sh}, bounded by a dense swept-up shell of interstellar material. The radius of the wind bubble can be several tens of parsecs. The very hot dilute bubble contains a small amount of mass, typically a few solar masses or even less, whereas the thin swept-up shell ultimately contains thousands of solar masses. This shell will have a gradual inner boundary where localized mixing of bubble and shell material proceeds. Since the typical SN ejecta mass is about $M_{ej} \approx 10 M_\odot$ for massive progenitors, only a relatively small fraction of the SN explosion energy E_{SN} is transformed into internal gas and CR energy during the SN shock propagation through the bubble and the main part of SN energy is deposited in the shell.

There are no physical reasons for which the CR production process in the shell matter is substantially different from that in the uniform ISM case, apart from the different distribution of ion injection across the SNR shock surface due to the magnetic field geometry (Völk et al. 2003). The major difference which plays a role for the π^0-decay γ-rays is that CRs accelerated during the SN shock propagation through the shell are confined in a medium which is more than a factor of ten denser than the average ISM in the Galactic disk, of density N_g^{GCR}. This may lead to an appreciable increase of the overall γ-ray production.

In addition, the SNR shock evolves differently than in the case of a uniform circumstellar medium. The entire active period of SNR evolution, when the SN shock is strong enough and effectively produces CRs, takes place within the shell of thickness L which is much smaller than its radius, $L \ll R_{sh}$. This implies that during the entire active evolution the SNR shock size R_s remains nearly constant. The time dependence of the shock speed V_s which is very important for the SCR confinement in SNRs can then be estimated as follows.

The majority of the progenitor stars with an intense wind are main-sequence stars with initial masses $M_i > 15 M_\odot$ (Abbott 1982). In the mean, during their evolution in the surrounding uniform ISM of gas number density $\rho_0 = m_p N_g^{GCR}$, they create a bubble of size (Weaver et al. 1977; Chevalier & Liang 1989)

$$R_{sh} = 0.76(0.5\dot{M}V_w^2 t_w^3/\rho_0)^{1/5}, \tag{20}$$

where $\dot{M}$ is the mass-loss rate of the presupernova star, V_w is the wind speed, and t_w is the duration of the wind period.

In order to determine the SN shock dynamics inside the shell we model the gas number density distribution in the bubble and in the shell in the form

$$N_g = N_b + (r/R_{sh})^{3(\sigma-1)} N_{sh}, \tag{21}$$

where $N_{sh} = \sigma N_g^{GCR}$ is the peak number density in the shell, N_b is the gas number density inside the bubble, typically very small compared with the shell density, and $\sigma = N_{sh}/N_0$ is the shell compression ratio.

The mass of the bubble

$$M_b = (4\pi R_{sh}^3/3) m_p N_b \tag{22}$$

is about 3 t*o* $5 M_\odot$, whereas the shell mass

$$M_{sh} = 4\pi N_{sh} m_p \int_0^{R_{sh}} dr r^2 (r/R_{sh})^{3(\sigma-1)} = (4\pi R_{sh}^3/3) N_g^{GCR} m_p \tag{23}$$

is a few thousands of solar masses. Therefore, during SNR shock propagation through the bubble, only a small fraction of its energy is given to gas of stellar origin. The main part of the explosion energy is deposited in the shell.

As in the case of a uniform ISM the evolution of the SN shock can be represented as a sequence of two stages.

The first stage is the free expansion phase which continues up to the time, when the shock sweeps up a mass equal to the ejecta mass M_{ej}. The second phase is analogous to the Sedov phase in the uniform ISM. The only significant difference compared with the uniform ISM case is that at the end of the free expansion phase, when a large part of the SN energy has been given to the swept up gas, the shock volume is much larger.

When the swept-up mass becomes large compared to the sum of bubble mass and M_{ej}, the SN shock propagation in the shell medium with density $N_g \propto r^{3(\sigma-1)}$ can be shown to approach the following self-similar adiabatic solution

$$R_s^3 V_s^2 N_g(R_s) = \text{const}, \tag{24}$$

which corresponds to an expansion law $R_s = R_0(t/t_0)^\nu$ with $\nu = 2/(3\sigma+2)$. The shock decelerates much more rapidly in this case, $V_s \propto t^{-3\sigma/(3\sigma+2)}$, than in the case of a uniform ISM, where $V_s \propto t^{-3/5}$. Due to the fact that $3\sigma \gg 1$ we have approximately $V_s \propto t^{-1}$. During this period of time the shock size and the amount of the swept-up mass are connected by the relation

$$R_s = R_{sh}(M/M_{sh})^{1/3\sigma}. \tag{25}$$

Substituting in this expression the SN ejecta mass M_{ej} we get the value of the SN shock size which corresponds to the end of free expansion phase

$$R_s(t_0) = R_0 = R_{sh}(M_{sh}/M_{ej})^{-1/3\sigma}. \tag{26}$$

At this epoch the shock speed is roughly the mean ejecta speed

$$V_s \approx V_0 = \sqrt{2E_{SN}/M_{ej}}. \tag{27}$$

Therefore the corresponding time scale is $t_0 = R_0/V_0$.

As it follows from the expansion law and from eq.(24), the SNR shock reaches the edge of the shell $R_s(t_f) = R_{sh}$ at time

$$t_f = t_0(M_{sh}/M_{ej})^{(3\sigma+2)/6\sigma}, \tag{28}$$

which for $\sigma \gg 1$ gives

$$t_f \approx t_0\sqrt{M_{sh}/M_{ej}}. \tag{29}$$

At about this stage the SNR shock has lost much of its speed and will come into pressure equilibrium with its environment, i.e. $t_f \approx T_{SN}$.

For the typical values $\dot{M} = 6 \times 10^{-9} M_\odot$/yr, $V_w = 2500$ km/s, $t_w = 10^7$ yr, which correspond to a progenitor of mass $M_i = 15M_\odot$ (Abbott 1982; Chevalier & Liang 1989) and $N_g^{GCR} = 1$ cm^{-3}, we have $R_{sh} \approx 30$ pc and $M_{sh} = 3 \times 10^3 M_\odot$. For $E_{SN} = 2 \times 10^{51}$ erg and $M_{ej} = 5M_\odot$ this gives $t_0 = 5 \times 10^3$ yr and $t_f \approx 10^5$ yr. The ambient interstellar gas may also be photoionized which implies a large enough external pressure to stop bubble expansion at a lower radius (Chevalier & Liang 1989). Therefore we also consider below the case $R_{sh} = 17$ pc which corresponds to $t_f = 3 \times 10^4$ yr.

According to theory, the expanding SNR shock produces a power law CR spectrum up to a maximum energy (Berezhko 1996; Berezhko et al. 1996; Berezhko & Völk 1997)

$$\epsilon_m \propto R_s V_s, \tag{30}$$

which is determined by the radius R_s and speed V_s of the shock. The CRs with the highest energy ϵ_{max} are produced at the very beginning of the Sedov phase $t \sim t_0$ when the product $R_s V_s$ has its maximum. Subsequently, the product $R_s V_s$ decreases with time approximately as t^{-1} and the SNR shock produces CRs with progressively lower cutoff energy $\epsilon_m(t) < \epsilon_{max} = \epsilon_m(t_0)$. During that phase those CRs that were previously produced with energies $\epsilon_m < \epsilon < \epsilon_{max}$ now leave the remnant without a significant influence of the SNR shock. Therefore the CR confinement time has to be taken in the form

$$T_p = min\{1, (t_0/t_f)(\epsilon_{max}/\epsilon)\} T_{SN}, \tag{31}$$

where according to the above consideration $T_{SN} = t_f$. This relation is analogous to eq.(3) in the case of a uniform ISM and leads to an increase of the spectral index of the γ-ray emission by one unit.

In order to estimate the highest energy of accelerated CRs in this case we use the relation $\epsilon_{max} \propto R_0 V_0 B$. Compared with the case of a uniform ISM, where $R_0 = 4$ pc for $N_g^{GCR} = 1$ cm^{-3}, the value of R_0 is about ten times larger but V_0 is the same. The magnetic field $B(t_0)$ on the other hand is much lower than in the case of a SN Ia. Indeed, as follows from eqs. (25) and (20), the gas number density at the beginning of the Sedov phase is as small as $N_g(t_0) = N_{sh} M_{ej}/M_{sh} \approx N_g^{GCR}/300$. The magnetic field in the region of the shell is presumably adiabatically compressed interstellar field B_0. If we approximate the directions of B_0 as being isotropic, then its magnitude scales with the gas density like $B \propto N_g^{2/3}$, e.g. Chevalier (1974). This gives a field strength $B(t_0) \approx 0.02 B_0$. Taking into account all the factors considered, we conclude that the maximum CR energy in the case of wind-SNe with dense shells is roughly a factor of ten smaller than in the case of SNe expanding into a uniform ISM. Assuming that the field amplification is proportional to $\sqrt{N_g}$, cf. Bell & Lucek (2001), this amplification does not play a role here since $N_g(t_0)$ is so low.

In contrast to our previous study (Berezhko & Völk 2000b), where we considered the CR acceleration by SN shocks expanding into a *modified* bubble, we consider here the opposite extreme of bubble structure. It is assumed that the magnetic field suppresses the mass and heat transport between the dense shell and the hot bubble (Chevalier & Liang 1989). Therefore the bubble mass is so small that a significant number of CRs is produced by the SN shock only when it enters the dense shell. Even though it is not clear at the moment which of these two concepts is physically more correct, we consider here the possibility of an unmodified bubble, because of its great importance for γ-ray production.

We note that due to the energy dependence of the confinement time $T_p(\epsilon)$ the SCR contribution to the γ-ray flux undergoes a spectral break at a γ-ray energy which corresponds to the SCR energy

$$\epsilon_{break} = (t_0/t_f)\epsilon_{max}. \tag{32}$$

Protons with energies $\epsilon \leq \epsilon_{break}$ survive up to the final active SNR epoch $t = t_f$ and therefore provide the maximum contribution to the γ-ray SNR emissivity at $\epsilon_\gamma^{break} = 0.1\epsilon_{break}$, whereas for $\epsilon_\gamma > \epsilon_\gamma^{break}$ the γ-ray emissivity decreases more steeply, to go to zero at $\epsilon_\gamma = 0.1\epsilon_{max}$.

The gas number density N_g^{SCR} in expression (2) is a function of SNR age t for wind-SNe, because the gas density $N_g(R_s)$ seen by the SN shock changes during the evolution of the SNR. It then follows from eq. (20) that the gas density behind the SN shock, where most of the SCRs are located, can be represented as $N_2(t) = \sigma_{SN} N_g(R_s(t)) = \sigma_{SN}(t/t_f)^2 N_{sh}$, where σ_{SN} is the SN shock compression ratio. Therefore the gas density seen by SCRs of energy ϵ is

$$N_g^{SCR} = \sigma_{SN}(T_p/T_{SN})^2 N_{sh}. \quad (33)$$

Since $T_p \propto \epsilon_\gamma^{-1}$ for $\epsilon_\gamma > \epsilon_\gamma^{break}$, the γ-ray spectrum produced in the shell is very steep at these high energies, and in eq. (2) we have $R(\epsilon_\gamma) \propto \epsilon_\gamma^{-2.4}$.

We assume that the progenitors of type II SNe are stars more massive than $M_i > 8M_\odot$, and that only those with $M_i > 15M_\odot$ have strong winds which produce extended bubbles, e.g. Abbott (1982). This is also true for the progenitor population of type Ib SNe. According to Güsten & Mezger (1983) the initial mass function has the form $dN/dM_i \propto M_i^{-\alpha}$ with $\alpha = 1.6, 2.4, 3.24, 3.62$ for the mass intervals $M_i < 1M_\odot$, $1M_\odot < M_i < 10M_\odot$, $10M_\odot < M_i < 50M_\odot$, $50M_\odot < M_i$, respectively. Therefore the ratio of the number of stars whose initial mass M_i exceeds 15 $M_\odot$ to those with $M_i > 8M_\odot$ is 0.23. Since about 85% of all SN explosions in the Galaxy are type II and type Ib SNe (Tammann et al. 1994), and since in addition the explosion energies E_{SN} of type II and type Ib SNe with $M_i > 15M_\odot$ are about two times larger than on average (Chevalier 1977; Hamuy 2003), we conclude that roughly a fraction $\delta = 0.3$ of the total Galactic SN energy release is from wind-SNe with progenitors of masses $M_i > 15M_\odot$ which produce extended bubbles.

Therefore the expected γ-ray SNR luminosity should be weighted between the SNR populations expanding into a uniform ISM and into bubbles as follows:

$$R = (1 - \delta)R_I + \delta R_{II}, \quad (34)$$

where R_I corresponds to the ratio of γ-ray production rates due to SCRs and GCRs if all SNe were exploding into a uniform ISM, and R_{II} is the corresponding ratio for explosions into bubbles.

Assuming that wind-SNe have rarefied bubbles with dense shells the spectrum of γ-rays calculated with the above ratio R, $\sigma = 10$, $\sigma_{SN} = 5$ and $\delta = 0.3$ is presented in Figs.1 and 2 by the full lines, which correspond to two different assumptions about the value of the SCR confinement time T_{SN}. It is seen that the wind-SNe dominate at energies $\epsilon_\gamma < 100$ GeV and fit the EGRET data for $T_{SN} = 3 \times 10^4$ yr fairly well, whereas for the larger confinement

times $T_{SN} = 10^5$ yr their contribution becomes so large, that their calculated flux exceeds the EGRET data by about a factor of two.

We emphasize that the contribution from the population of wind-SNe is mainly due to the nuclear SCR component. The break in the γ-ray spectra presented in Fig.1, 2 is at $\epsilon_\gamma^{break} = 30$ GeV for $T_{SN} = 10^5$ yr and at $\epsilon_\gamma^{break} = 100$ GeV for $T_{SN} = 3 \times 10^4$ yr. For larger energies the proton confinement time and the gas density become lower, $T_p \propto \epsilon^{-1}$, $N_g^{SCR} \propto \epsilon^{-2}$, which leads to a corresponding decrease of the wind-SN contribution to the γ-ray spectrum. This contribution becomes insignificant for $\epsilon_\gamma \sim 1$ TeV.

4. Discussion

Our above consideration is based upon a picture in which the population of SNRs, considered here as the main GCR sources, are distributed across the Galactic disk similar to the Interstellar gas. In this case the calculated ratio $R(\epsilon_\gamma)$ of the SCR to the GCR contribution does not depend upon the observing direction in the disk. Since the number of these sources $N_{SN} = \nu_{SN} T_{SN}$ is very limited, their mean number within the field of view $\sim$ few degrees of a stereoscopic system of imaging atmospheric Cherenkov telescopes for γ-ray energies $\gtrsim$100 GeV is so low, that one should expect large fluctuations of the actual value of $R(\epsilon_\gamma)$, especially for lines-of-sight other than those directed towards the inner Galaxy.

The actual SNR distribution within the Galactic disk most probably is not uniform. According to Case & Bhattacharya (1998) the SN explosion rate as a function of Galactocentric radius r has a peak at $r \approx 5$ kpc and drops exponentially with a scale length of ~ 7 kpc. Within $5 < r < 20$ kpc this agrees fairly well with the radial distributions of *Supernovae* in a sample of 36 *external galaxies* of Hubble type Sb–Sbc, statistically considered equivalent to our Galaxy (Dragicevich et al. 1999). From our results, for γ-ray energies larger than 10 GeV, the ratio R, corrected for the nonuniform SNR distribution, is then given by the ratio $R_c = \eta(r)R$, where the parameter $\eta(r)$ represents the ratio of the actual SNR number along the given line-of-sight to that SNR number which would corresponds to their uniform distribution. For lower γ-ray energies, $100 < \epsilon_\gamma < 10^4$ MeV, where the truly diffuse emission from the GCRs dominates, it is known that the γ-ray emissivity gradient is significantly shallower than the gradient of $\eta(r)$, as shown by the EGRET instrument on board of the CGRO satellite (Strong & Mattox 1996) in basic agreement with the results from the COS-B satellite (Strong et al. 1988). This much shallower, indeed truly diffuse Galactocentric γ-ray gradient can be understood as a nonlinear propagation effect from the disk into the Galactic Wind where the latter is driven by the GCRs themselves (Breitschwerdt et al. 2002). In contrast,

the radial galactocentric variation of the direct radiation from the sources is not smoothed by any propagation effect. Therefore its amplitude should spatially vary like the distribution of the sources, integrated along the line-of-sight, and this variation should become directly measurable in the > 100 GeV range. Given several such line-of-sight integrals, one can infer the radial galactocentric γ-ray emissivity gradient. The excellent angular resolution of atmospheric Cherenkov telescopes and the high sensitivity of instruments like the CANGAROO III and H.E.S.S. arrays in the Southern Hemisphere, which at TeV energies have an order of magnitude higher sensitivity compared with previous generation instruments like HEGRA (Aharonian et al. 2001), should allow this measurement in particular in the inner Galaxy.

The assumed average value $B = 30$ μG of the magnetic field inside SNRs is considerably smaller than the value 120 μG inferred for e.g. SN 1006 in the very early Sedov phase (Berezhko et al. 2002, 2003a). However the amplification must be a monotonically increasing function of the shock strength. In particular, Bell & Lucek (2001) argue for a dependence $B \propto V_s$. If so, then the effective amplified magnetic field in SNRs varies from $B \sim 100$ μG at the beginning of the Sedov phase to the typical interstellar value $B \approx 10$ μG at the end of its active period $t = T_{SN}$. Since the value of the magnetic field in SNRs has a dominant influence on the IC γ-ray spectrum, and since all evolutionary phases of a SNR contribute to this spectrum in a roughly equal way, the adopted value $B = 30$ μG appears to be quite realistic, averaging over the entire active period.

One should also note that the SNR magnetic field influences the maximum energy of SCRs. Since the most energetic SCRs are created at the beginning of the Sedov phase, their actual maximum energy $\epsilon_{max} \propto B(t_0)$ is expected to be larger than considered here, because the amplified magnetic field $B(t_0)$ in this phase is by a factor of several larger than on average. Compared with the Fig.2 this effect will produce more a extended and smooth high energy tail of the γ-ray spectrum up to the energy $\epsilon_\gamma \approx 10^5$ GeV.

The actual process of SN shock interaction with the dense shell could be much more complicated compared with the ideal picture considered here. If due to some physical factors (inhomogeneities in the surrounding ISM, instabilities ...) the shell is *always* strongly distorted, then one would expect that the SN shock will penetrate through the shell more rapidly. This will lead to a decrease of the propagation time t_f and/or to the decrease of the effective gas number density N_g^{SCR}, that in turn will decrease the amount of γ-rays produced in such types of SNRs. In other words, the calculated γ-ray spectrum at $\epsilon_\gamma \lesssim 100$ GeV (see Fig.1, 2), which is mainly due to the wind SNe, has to be considered as an upper limit, which corresponds to the case of stable unmodified bubbles.

5. Summary

Our considerations demonstrate that the SCRs inevitably make a strong contribution to the "diffuse" γ-ray flux from the Galactic disk at all energies above a few GeV, if the population of SNRs is the main source of the GCRs. According to our estimates, the SCR contribution dominates at energies greater than 100 GeV due to its substantially harder spectrum.

There are two physical parameters which influence the expected γ-ray emission from SNRs significantly: the CR confinement time and the mean magnetic field strength B inside the SNRs. For a conventional value $B = 10$ μG the expected γ-ray flux from SNRs exceeds the HEGRA upper limit considerably if the SCR confinement time is as large as $T_{SN} = 10^5$ yr. This contradiction can be resolved either if we suggest an appreciably higher postshock magnetic field $B \gtrsim 30$ μG or if the SCR confinement time is as small as $T_{SN} \sim 10^4$ yr, or a combination of both effects. These possibilities can be attributed to field amplification by the SCRs themselves. In fact, nonlinear field amplification may also lead to a substantial decrease of the SCR confinement time: according to Ptuskin & Zirakashvili (2003) maximal turbulent Alfvén wave damping with its corresponding increase of CR mobility could make the SCR confinement time as small as $T_{SN} \sim 10^4$ yr. Under these circumstances the most realistic γ-ray background spectrum is represented by the thick solid line in Fig.2. One can see that even in the case of $B = 30$ μG and $T_{SN} \sim 3 \times 10^4$ yr the SNR contribution at TeV energies, with roughly numbers of IC and π^0-decay γ-rays ; still exceeds the GCR contribution almost by an order of magnitude. Note that the preliminary γ-ray flux at $\epsilon_\gamma = 1$ TeV, measured by the Milagro detector (Fleysher 2003), confirms our earlier prediction (Berezhko & Völk 2003).

At lower energies $\epsilon_\gamma \lesssim 100$ GeV the γ-ray emission from SNRs is dominated by the wind-SNe with initial progenitor mass $M_i > 15 M_\odot$ which expand into the bubble created by the progenitor's wind, under the assumption that the bubble is not modified by global heat and mass transport. The main fraction of the CRs is produced in this case when the SN shock propagates through the thin dense shell at the edge of the bubble. Due to the significantly higher mass density of the shell material compared with the mean gas density in the Galactic disk π^0-decay γ-rays dominate at γ-ray energies $\epsilon_\gamma \lesssim 100$ GeV despite the fact that only about 20% of all SNRs belong to this class of objects. As shown in Figs.1 and 2, the discrepancy between the observed "diffuse" intensity and standard model predictions at energies above a few GeV can be attributed to the SCR contribution. This requires unmodified bubbles for the wind-SNe. Since we cannot prove this assumption, the corresponding explanation of the lower energy γ-ray excess is not a definitive conclusion, but rather a plausible suggestion which must await more detailed studies of wind bubble morphologies.

We conclude that a detailed measurement of the low-latitude “diffuse” Galactic γ-ray *spectrum* within the energy interval from 1 to 10^4 GeV will allow a strong consistency check for the predominant origin of the Galactic Cosmic Rays from the Galactic population of SNRs as a whole. In the energy range above a few 100 GeV our predictions are quite robust resulting in a hard spectrum whose Galactocentric variation should correspond to that of the observed SNR distribution. If the preliminary Milagro data, which agree with our prediction, are confirmed by more precise measurements, then such measurements will give an indirect confirmation of SNRs as the main sources of GCRs. The measurements will also give a new tool to study the spatial distribution of SNRs in the Galaxy.

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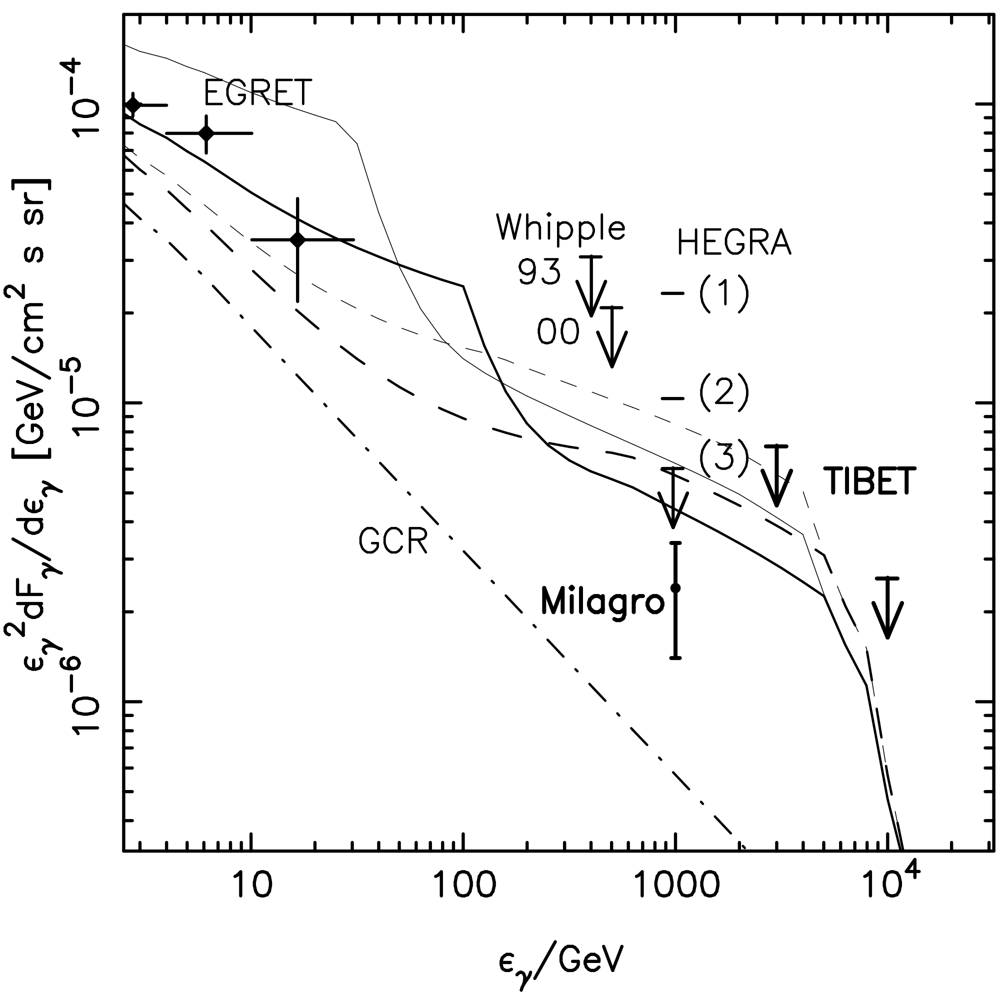

Fig. 1.— The average diffuse γ-ray spectrum of the low latitude inner Galaxy ($38^\circ < l < 43^\circ$, $|b| \leq 2^\circ$). The thick and the thin dashed lines represent the case of a uniform ISM with the two different SNR life times $T_{SN} = 10^5$ yr and $T_{SN} = 3 \times 10^4$ yr, respectively; the full lines correspond to the case when 30% of the SN energy release is due to SNRs which expand into the wind bubble of a massive progenitor star. EGRET data (Hunter et al. 1997b), preliminary Milagro data at $\epsilon_\gamma = 1$ TeV (Fleysher 2003), the Wipple upper limits (Reynolds et al. 1993; LeBohec et al. 2000), the HEGRA upper limits (Aharonian et al. 2001), and the Tibet upper limits (Amenomori et al. 2003) at $\epsilon_\gamma = 3$, 10 TeV are shown. The Milagro data have been multiplied by a factor of 3 (see text).

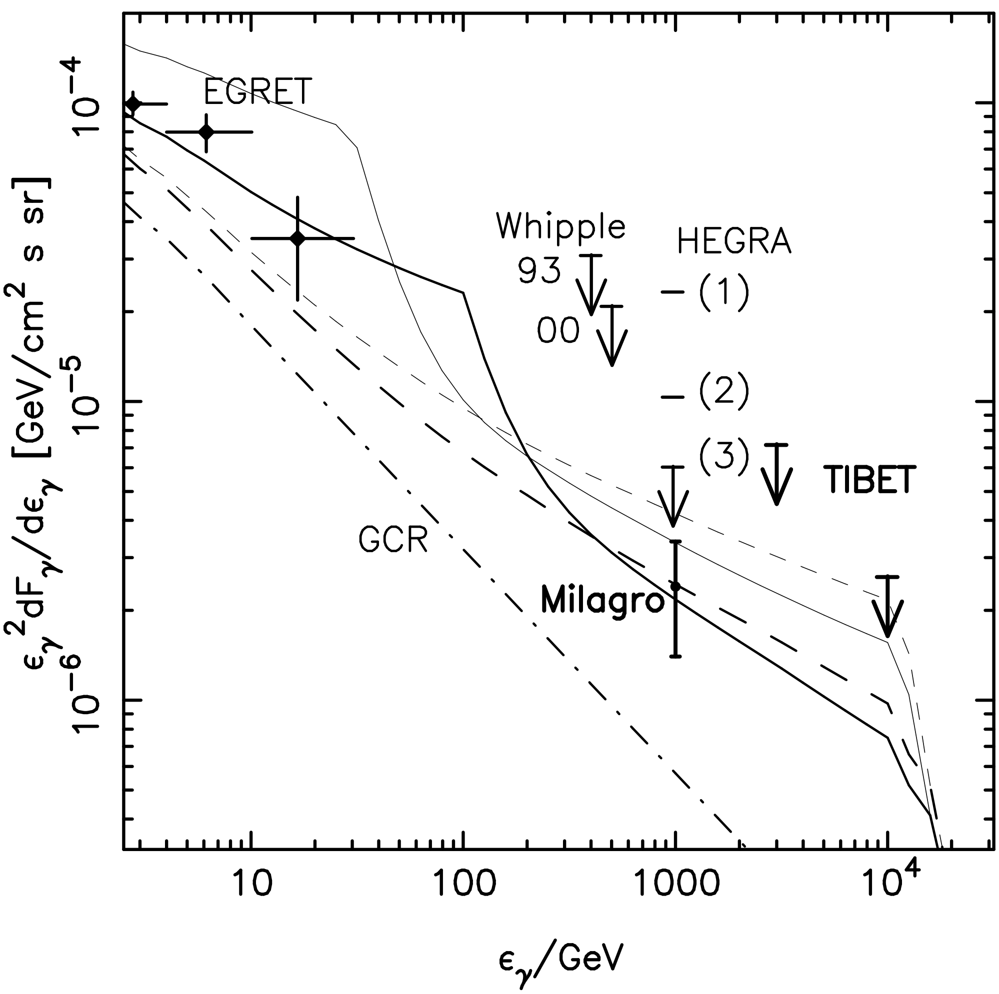

Fig. 2.— The same as in Fig.1 but with the postshock SNR magnetic field value $B = 30\ \mu$G.